

Electroweak chiral Lagrangians and the Higgs properties at the one-loop level

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Abstract. In these proceedings we explore the use of (non-linear) electroweak chiral Lagrangians for the description of possible beyond the Standard Model (BSM) strong dynamics in the electroweak (EW) sector. Experimentally one observes an approximate EW symmetry breaking pattern $SU(2)_L \times SU(2)_R / SU(2)_{L+R}$. Quantum Chromodynamics (QCD) shows a similar chiral structure [1] and, in spite of the differences (in the EW theory $SU(2)_L \times U(1)_Y$ is gauged), it has served for years as a guide for this type of studies [2–4]. Examples of one-loop computations in the low-energy effective theory and the theory including the first vector (V) and axial-vector (A) resonances are provided, yielding, respectively, predictions for $\gamma\gamma \rightarrow Z_L Z_L, W_L^+ W_L^-$ and the oblique parameters S and T .

1 Introduction: strong dynamics and chiral Lagrangians

A non-linear realization of the EW would-be Goldstone bosons (WBGBs) is considered to build the EW low-energy effective field theory (EFT), which is described by an EW chiral Lagrangian with a light Higgs (ECLh). It includes the Standard Model (SM) content: the EW Goldstones w^a , the EW gauge bosons W_μ^a and B_μ and a singlet Higgs h (the fermion sector is not discussed here). In particular, in Sec. 2 we explain the chiral counting in the ECLh [5, 6] and provide an example of a next-to-leading order (NLO) computation: we calculate $\gamma\gamma \rightarrow W_L^+ W_L^-, Z_L Z_L$ within this framework up to the one-loop level [5] at energies below new possible composite resonances, $\sqrt{s} \ll \Lambda_{\text{ECLh}} \sim \min\{M_R, 4\pi v\}$ (with $v = (\sqrt{2}G_F)^{-1/2}$ and $4\pi v \simeq 3$ TeV). Analogous works on WW -scattering can be found in Refs. [7].

However, in the case of having heavy composite resonance, the EFT stops being valid when the energy becomes of the order of their masses (expected to be of the order of $M_R \sim 4\pi v \sim 3$ TeV). One has to introduce these new degrees of freedom in our EW Lagrangian following a procedure analogous to that in QCD [8]. Likewise, under reasonable ultraviolet (UV) completion hypotheses like, e.g., the Weinberg sum-rules (WSRs) fulfilled by certain types of theories [9–13], one can make predictions on low-energy observables. In Sec. 3 we write down the relevant $SU(2)_L \times SU(2)_R$ invariant Lagrangian including the SM content and a multiplet of V and A resonances and extract one-loop limits on the resonance masses and the Higgs coupling g_{hWW} [14] from the experimental values of oblique parameters S and T [15]. Alternative one-loop analyses can be found in Refs. [10, 16].

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2 Low-energy EFT: ECLh and one-loop $\gamma\gamma \rightarrow W_L^a W_L^b$ scattering

The Higgs boson does not enter in the SM at tree-level in these processes (where one also has $\mathcal{M}(\gamma\gamma \rightarrow ZZ)_{\text{SM}}^{\text{tree}} = 0$). Nevertheless, one can search for new physics by studying the one-loop corrections [5], which are sensitive to deviations from the SM in the Higgs boson couplings. Our analysis [5] is performed in the Landau gauge and making use of the Equivalence Theorem (Eq.Th.) [17],

$$\mathcal{M}(\gamma\gamma \rightarrow W_L^a W_L^b) \simeq -\mathcal{M}(\gamma\gamma \rightarrow w^a w^b), \quad (1)$$

valid in the energy regime $m_W^2, m_Z^2 \ll s$. The EW gauge boson masses $m_{W,Z}$ are then neglected in our computation. Furthermore, since $m_h \sim m_{W,Z} \ll 4\pi v \simeq 3 \text{ TeV}$ we also neglect m_h in our calculation. In summary, the applicability range in [5] is

$$m_W^2, m_Z^2, m_h^2 \stackrel{\text{Eq.Th.}}{\ll} s, t, u \stackrel{\text{EFT}}{\ll} \Lambda_{\text{ECLh}}^2, \quad (2)$$

with the upper limit given by the EFT cut-off Λ_{ECLh} , expected to be of the order of $4\pi v \simeq 3 \text{ TeV}$ or the mass of possible heavy BSM particles.

The WBGBs are described by a matrix field U that takes values in the $SU(2)_L \times SU(2)_R / SU(2)_{L+R}$ coset, and transforms as $U \rightarrow LUR^\dagger$ [2, 3]. The relevant ECLh with the basic building blocks is

$$U = u^2 = 1 + iw^a \tau^a / v + \mathcal{O}(w^2), \quad D_\mu U = \partial_\mu U + i\hat{W}_\mu U - iU \hat{B}_\mu, \quad V_\mu = (D_\mu U)U^\dagger, \quad u^\mu = -iu^\dagger D^\mu U u^\dagger, \\ \hat{W}_{\mu\nu} = \partial_\mu \hat{W}_\nu - \partial_\nu \hat{W}_\mu + i[\hat{W}_\mu, \hat{W}_\nu], \quad \hat{B}_{\mu\nu} = \partial_\mu \hat{B}_\nu - \partial_\nu \hat{B}_\mu, \quad \hat{W}_\mu = gW_\mu^a \tau^a / 2, \quad \hat{B}_\mu = g' B_\mu \tau^3 / 2, \quad (3)$$

with well-defined transformation properties [3, 5, 14]. Two particular parametrizations of the unitary matrix U (exponential and spherical) were considered in [5], both leading to the same predictions for the physical (on-shell) observables.¹ We consider the counting $\partial_\mu, m_W, m_Z, m_h \sim \mathcal{O}(p)$, $D_\mu U, V_\mu \sim \mathcal{O}(p)$ and $\hat{W}_{\mu\nu}, \hat{B}_{\mu\nu} \sim \mathcal{O}(p^2)$ [5, 6]. We require the ECLh Lagrangian to be CP invariant, Lorentz invariant and $SU(2)_L \times U(1)_Y$ gauge invariant. Here we focus ourselves on the relevant terms for $\gamma\gamma \rightarrow w^a w^b$ at leading order (LO) $-\mathcal{O}(p^2)$ – and NLO in the chiral counting $-\mathcal{O}(p^4)$ – [3, 5]:

$$\mathcal{L}_2 = -\frac{1}{2g^2} \langle \hat{W}_{\mu\nu} \hat{W}^{\mu\nu} \rangle - \frac{1}{2g'^2} \langle \hat{B}_{\mu\nu} \hat{B}^{\mu\nu} \rangle + \frac{v^2}{4} \left[1 + 2a \frac{h}{v} + b \frac{h^2}{v^2} \right] \langle D^\mu U^\dagger D_\mu U \rangle + \frac{1}{2} \partial^\mu h \partial_\mu h + \dots, \quad (4)$$

$$\mathcal{L}_4 = a_1 \text{Tr}(U \hat{B}_{\mu\nu} U^\dagger \hat{W}^{\mu\nu}) + ia_2 \text{Tr}(U \hat{B}_{\mu\nu} U^\dagger [V^\mu, V^\nu]) - ia_3 \text{Tr}(\hat{W}_{\mu\nu} [V^\mu, V^\nu]) - \frac{c_\gamma}{2} \frac{h}{v} e^2 A_{\mu\nu} A^{\mu\nu} + \dots$$

where $\langle X \rangle$ stands for the trace of the 2×2 matrix X , one has the photon field strength $A_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ and the dots stand for operators not relevant within our approximations for $\gamma\gamma$ -scattering [5].

The amplitudes $\mathcal{M}(\gamma(k_1, \epsilon_1)\gamma(k_2, \epsilon_2) \rightarrow w^a(p_1)w^b(p_2))$, with $w^a w^b = zz, w^+ w^-$, have the structure

$$\mathcal{M} = ie^2 (\epsilon_1^\mu \epsilon_2^\nu T_{\mu\nu}^{(1)}) A(s, t, u) + ie^2 (\epsilon_1^\mu \epsilon_2^\nu T_{\mu\nu}^{(2)}) B(s, t, u), \quad (5)$$

written in terms of the two independent Lorentz structures $T_{\mu\nu}^{(1,2)} \sim \mathcal{O}(p^2)$ involving the external momenta, which can be found in [5]. The Mandelstam variables are defined as $s = (p_1 + p_2)^2$, $t = (k_1 - p_1)^2$ and $u = (k_1 - p_2)^2$ and the ϵ_i 's are the polarization vectors of the external photons.

In dimensional regularization, our NLO computation of the $\mathcal{M}(\gamma\gamma \rightarrow w^a w^b)$ amplitudes can be systematically sorted out in the form [5]

$$\mathcal{M} = \mathcal{M}_{\text{LO}} + \mathcal{M}_{\text{NLO}} \sim \underbrace{\mathcal{O}(e^2)}_{\text{LO, tree}} + \left[\underbrace{\mathcal{O}\left(e^2 \frac{p^2}{16\pi^2 v^2}\right)}_{\text{NLO, 1-loop}} + \underbrace{\mathcal{O}\left(e^2 \frac{a_i p^2}{v^2}\right)}_{\text{NLO, tree}} \right], \quad (6)$$

¹ Other representations have been recently studied in Ref. [18].

	ECLh	ECL (Higgsless)
$\Gamma_{a_1-a_2+a_3}$	0	0
Γ_{c_γ}	0	-
Γ_{a_1}	$-\frac{1}{6}(1-a^2)$	$-\frac{1}{6}$
$\Gamma_{a_2-a_3}$	$-\frac{1}{6}(1-a^2)$	$-\frac{1}{6}$
Γ_{a_4}	$\frac{1}{6}(1-a^2)^2$	$\frac{1}{6}$
Γ_{a_5}	$\frac{1}{8}(b-a^2)^2 + \frac{1}{12}(1-a^2)^2$	$\frac{1}{12}$

Table 1. Running $\frac{da_i^r}{d \ln \mu} = -\frac{\Gamma_{a_i}}{16\pi^2}$ of the relevant ECLh parameters and their combinations appearing in the six selected observables. The third column provides the corresponding running for the Higgsless EW chiral Lagrangian (ECL) case [4]. For the sake of completeness, we have added the running of the ECLh parameters a_4^r and a_5^r , which has been recently determined in the one-loop analysis of WW -scattering within the framework of chiral Lagrangians [7]. One can see that in the SM limit ($a = b = 1$) these \mathcal{L}_4 coefficients do not run, in agreement with the fact that these higher order operators are absent in the SM.

where $e \sim \mathcal{O}(p/v)$ and A and B are given up to NLO by [5]

$$\begin{aligned}
A(\gamma\gamma \rightarrow zz)_{\text{LO}} &= B(\gamma\gamma \rightarrow zz)_{\text{LO}} = 0, \\
A(\gamma\gamma \rightarrow zz)_{\text{NLO}} &= \frac{2ac_\gamma^r}{v^2} + \frac{(a^2-1)}{4\pi^2 v^2}, & B(\gamma\gamma \rightarrow zz)_{\text{NLO}} &= 0, \\
A(\gamma\gamma \rightarrow w^+w^-)_{\text{LO}} &= 2sB(\gamma\gamma \rightarrow w^+w^-)_{\text{LO}} = -\frac{1}{t} - \frac{1}{u}, \\
A(\gamma\gamma \rightarrow w^+w^-)_{\text{NLO}} &= \frac{2ac_\gamma^r}{v^2} + \frac{(a^2-1)}{8\pi^2 v^2} + \frac{8(a_1^r - a_2^r + a_3^r)}{v^2}, & B(\gamma\gamma \rightarrow w^+w^-)_{\text{NLO}} &= 0.
\end{aligned} \tag{7}$$

The term with c_γ^r comes from the Higgs tree-level exchange in the s -channel, the term proportional to (a^2-1) comes from the one-loop diagrams with \mathcal{L}_2 vertices, and the Higgsless operators in Eq. (5) yield the tree-level contribution to $\gamma\gamma \rightarrow w^+w^-$ proportional to $(a_1 - a_2 + a_3)$. Independent diagrams are in general UV divergent. However, in dimensional regularization, the final one-loop amplitude turns out to be UV finite and one has $a_1^r - a_2^r + a_3^r = a_1 - a_2 + a_3$, $c_\gamma^r = c_\gamma$ [5], as in the Higgsless case [19].

In order to pin down each of the relevant combinations of ECLh couplings in Eq. (7) (a , c_γ^r and $a_1^r - a_2^r + a_3^r$) one must combine our $\gamma\gamma$ -scattering analysis with other observables that depend on this same set of parameters. It is not difficult to find that other processes involving photons depend on these parameters. In Ref. [5] we computed 4 more observables of this kind: the $h \rightarrow \gamma\gamma$ decay width (depending on a and c_γ), the oblique S -parameter (depending on a and a_1), and the $\gamma^* \rightarrow w^+w^-$ (depending on a and $a_2 - a_3$) and $\gamma^*\gamma \rightarrow h$ (depending on c_γ) electromagnetic form-factors. The one-loop contribution in these six relevant amplitudes is found to be UV-divergent in some cases. These divergences are absorbed by means of the generic $\mathcal{O}(p^4)$ renormalizations $a_i^r(\mu) = a_i + \delta a_i$. As expected, the renormalization in the six observables gives a fully consistent set of renormalization conditions and fixes the running of the renormalized couplings in the way given in Table 1.

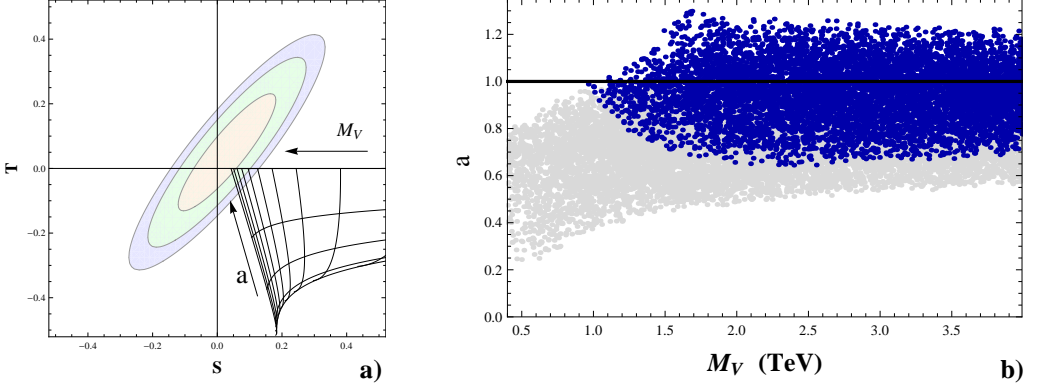


Figure 1. **a)** NLO determinations of S and T , imposing the two WSRs. The approximately vertical curves correspond to constant values of M_V , from 1.5 to 6.0 TeV at intervals of 0.5 TeV. The approximately horizontal curves have constant values of a : 0.00, 0.25, 0.50, 0.75, 1.00. The arrows indicate the directions of growing M_V and a . The ellipses give the experimentally allowed regions at 68% (orange), 95% (green) and 99% (blue) confidence level (CL). **b)** Scatter plot for the 68% CL region, in the case when only the first WSR is assumed. The dark blue and light gray regions correspond, respectively, to $0.2 < M_V/M_A < 1$ and $0.02 < M_V/M_A < 0.2$.

3 Impact of spin-1 composite resonances on the oblique parameters

One can extend the range of validity and predictability of the ECLh by adding possible new states to the theory. Thus, the lightest V and A resonances are added to the EW Lagrangian in Ref. [14] in order to describe the oblique parameters S and T [9]. The relevant EW chiral invariant Lagrangian is given by the kinetic and Yang-Mills terms and the interactions [14]^{2, 3}

$$\begin{aligned} \mathcal{L} = & \frac{v^2}{4} \langle u_\mu u^\mu \rangle \left(1 + \frac{2a}{v} h \right) + \frac{F_V}{2\sqrt{2}} \langle V_{\mu\nu} f_+^{\mu\nu} \rangle + \frac{iG_V}{2\sqrt{2}} \langle V_{\mu\nu} [u^\mu, u^\nu] \rangle \\ & + \frac{F_A}{2\sqrt{2}} \langle A_{\mu\nu} f_-^{\mu\nu} \rangle + \sqrt{2} \lambda_1^{hA} \partial_\mu h \langle A^{\mu\nu} u_\nu \rangle. \end{aligned} \quad (8)$$

In order to compute S and T up to the one-loop level we use the dispersive representations [9, 14],

$$S = \frac{16\pi}{g^2 \tan \theta_W} \int_0^\infty \frac{dt}{t} [\rho_S(t) - \rho_S(t)^{\text{SM}}], \quad T = \frac{4\pi}{g'^2 \cos^2 \theta_W} \int_0^\infty \frac{dt}{t^2} [\rho_T(t) - \rho_T(t)^{\text{SM}}], \quad (9)$$

with $\rho_S(t)$ the spectral function of the $W^3 B$ correlator [9, 21] and $\rho_T(t)$ the spectral function of the difference of the neutral and charged Goldstone self-energies [14]. The calculation of T above has been simplified by means of the Ward-Takahashi relation $T = Z^{(w^+)} / Z^{(w^0)} - 1$ [20]. Only the lightest two-particle cuts have been considered in $\rho_S(t)$ and $\rho_T(t)$, respectively, $\{w\bar{w}, w\bar{h}\}$ and $\{B\bar{w}, B\bar{h}\}$. Since $\rho_S(t)^{\text{SM}} \xrightarrow{t \rightarrow \infty} 0$, the convergence of the Peskin-Takeuchi sum-rule requires $\rho_S(t) \xrightarrow{t \rightarrow \infty} 0$. Furthermore, assuming that weak isospin and parity are good symmetries of the BSM strong dynamics, the $W^3 B$ correlator is proportional to the difference of the vector and axial-vector two-point Green's functions [9]. In asymptotically-free gauge theories this difference vanishes at $s \rightarrow \infty$ as $1/s^3$ [12], implying the (tree-level) LO WSRs [13],

$$F_V^2 - F_A^2 = v^2 \quad (\text{1st WSR}), \quad F_V^2 M_V^2 - F_A^2 M_A^2 = 0 \quad (\text{2nd WSR}). \quad (10)$$

² Here we follow the notation $f_\pm^{\mu\nu} = u^\dagger \hat{W}^{\mu\nu} u \pm u \hat{B}^{\mu\nu} u^\dagger$ from Ref. [14, 21], where there is a global sign difference with [5] in the definitions of \hat{W}_μ and \hat{B}_μ . The spin-1 resonances are described in the antisymmetric tensor formalism [8].

³ In other works, the coupling a can be found with the notation κ_W and ω [14] or κ_V [22].

	a	M_V
two WSRs	$0.97 - 1$	$> 5 \text{ TeV}$
Only 1st WSR: $0.2 < M_V/M_A < 1$	$0.6 - 1.3$	$> 1 \text{ TeV}$
$0.5 < M_V/M_A < 1$	$0.84 - 1.30$	$> 1.5 \text{ TeV}$
$M_V/M_A = 1$	$0.97 - 1.30$	$> 1.8 \text{ TeV}$
$(M_V > 1 \text{ TeV})^\dagger \quad 1 < M_V/M_A < 2$	$0.7 - 1.9$	$> 1 \text{ TeV}^\dagger$

Table 2. Allowed range for the M_V and a at the 68% CL for the two-WSRs (where V and A are very degenerate since $M_V^2/M_A^2 = a$ in this case) and only-1st-WSR cases (for various values M_V/M_A). In the last line we also impose the restriction[†] $M_V > 1 \text{ TeV}$.

However, although the 1st WSR is expected to be true in gauge theories with non-trivial ultraviolet fixed points [10, 11], the 2nd WSR is questionable in some of these models. Thus, two alternative scenarios are studied in Ref. [14]: one assuming the two WSRs and another assuming just the 1st WSR. At tree-level one has the LO determinations [9, 14, 21]

$$S_{\text{LO}} = 4\pi \left(\frac{F_V^2}{M_V^2} - \frac{F_A^2}{M_A^2} \right) = \frac{4\pi v^2}{M_V^2} \left(1 + \frac{M_V^2}{M_A^2} \right) \quad (\text{1st \& 2nd WSR}), \quad (11)$$

$$S_{\text{LO}} = 4\pi \left\{ \frac{v^2}{M_V^2} + F_A^2 \left(\frac{1}{M_V^2} - \frac{1}{M_A^2} \right) \right\} > \frac{4\pi v^2}{M_V^2} \quad (\text{1st WSR \& } M_V < M_A).$$

In the first case, the two WSRs imply $M_V < M_A$ and determine F_V and F_A in terms of the resonance masses [8, 9, 14, 21]. In the second case, it is not possible to extract a definite prediction with just the 1st WSR but one can still derive the inequality above if one assumes a similar mass hierarchy $M_V < M_A$. On the other hand, this inequality flips direction if $M_A < M_V$ or turns into an equality in the degenerate case $M_V = M_A$ [14]. At NLO the computed W^3B correlator is given by the ww and hw cuts, whose contributions to the $\rho_S(t)$ spectral function would have an unphysical grow at high energies unless $F_V G_V = v^2$ and $F_A \lambda_1^{hA} = av$ [8, 14, 21]. Thus, we obtain the NLO prediction [14]

$$S = 4\pi v^2 \left(\frac{1}{M_V^2} + \frac{1}{M_A^2} \right) \quad (12)$$

$$+ \frac{1}{12\pi} \left[\log \frac{M_V^2}{m_H^2} - \frac{11}{6} + \frac{M_V^2}{M_A^2} \log \frac{M_A^2}{M_V^2} - \frac{M_V^4}{M_A^4} \left(\log \frac{M_A^2}{m_{S_1}^2} - \frac{11}{6} \right) \right] \quad (\text{1st \& 2nd WSR}),$$

$$S > \frac{4\pi v^2}{M_V^2} + \frac{1}{12\pi} \left[\log \frac{M_V^2}{m_H^2} - \frac{11}{6} - a^2 \left(\log \frac{M_A^2}{m_{S_1}^2} - \frac{17}{6} + \frac{M_A^2}{M_V^2} \right) \right], \quad (\text{1st WSR \& } M_V < M_A).$$

In the two-WSRs scenario, in order to enforce the 2nd WSR at NLO one needs the additional constraint $a = M_V^2/M_A^2$ (hence restricted to the range $0 \leq a \leq 1$). Again, the inequality in the last line flips direction or turns into an equality when, respectively, $M_A < M_V$ or $M_V = M_A$.

At LO, $\rho_T(t)$ is zero and one has $T_{\text{LO}} = 0$. At NLO, where we enforce the $\rho_S(t)$ constraints $F_V G_V = v^2$ and $F_A \lambda_1^{hA} = av$, we find that $\rho_T(t) \xrightarrow{t \rightarrow \infty} 0$ and obtain the NLO prediction

$$T = \frac{3}{16\pi \cos^2 \theta_W} \left[1 + \log \frac{m_h^2}{M_V^2} - a^2 \left(1 + \log \frac{m_h^2}{M_A^2} \right) \right], \quad (13)$$

In Fig. 1, we show the compatibility between the experimental determinations for S and T [15] and our NLO determinations in both scenarios. The numerical results in Table 2 show that the precision

electroweak data requires resonance masses over the TeV and the hWW coupling to be close to the SM one ($a^{\text{SM}} = 1$), in agreement with present LHC bounds [22].

To conclude, we emphasize that, remarkably, just by considering the experimental m_h (the only LHC input) and the EW precision observables (LEP input), the allowed region concentrates around $a \simeq 1$ for reasonable values of the splitting $M_V/M_A \sim \mathcal{O}(1)$ (see Fig. 1 and Table 2).

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